

Super-Eddington accretion in ultra-luminous neutron star binary

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ABSTRACT

We discuss properties of the ultra-luminous X -ray source in the galaxy M82, NuSTAR J095551+6940.8, containing an accreting neutron star. The neutron star has surface magnetic field $B_{NS} \approx 1.4 \times 10^{13}$ G and experiences accretion rate of $9 \times 10^{-7} M_{\odot} \text{ yr}^{-1}$. The magnetospheric radius, close to the corotation radius, is $\sim 2 \times 10^8$ cm. The accretion torque on the neutron star is reduced well below what is expected in a simple magnetospheric accretion due to effective penetration of the stellar magnetic field into the disk beyond the corotation radius. As a result, the radiative force of the surface emission does not lead to strong coronal wind, but pushes plasma along magnetic field lines towards the equatorial disk. The neutron star is nearly an orthogonal rotator, with the angle between the rotation axis and the magnetic moment ≥ 80 degrees. Accretion occurs through optically thick – geometrically thin and flat accretion “curtain”, which cuts across the polar cap. High radiation pressure from the neutron star surface is nevertheless smaller than that the ram pressure of the accreting material flowing through the curtain, and thus fails to stop the accretion. At distances below few stellar radii the magnetic suppression of the scattering cross-section becomes important. The X -ray luminosity (pulsed and persistent components) comes both from the neutron star surface as a hard X -ray component and as a soft component from reprocessing by the accretion disk.

1. Introduction

Ultra-luminous X -ray sources (ULXs) is a class of objects observed in nearby galaxies with luminosities in the $10^{39} - 10^{41} \text{ erg s}^{-1}$ range, not associated with the central black hole (Fabbiano & White (2006); Roberts (2007); Feng & Soria (2011)). Such luminosities exceed Eddington luminosity of a stellar-mass central object. Until recently, the discussion on the nature of these objects concentrated on a choice between jetted/beamed emission from stellar mass black holes or accretion onto intermediate, $M \sim 100 - 1000 M_{\odot}$, black hole (possibly also producing a jet) (Kaaret et al. 2003; Mapelli et al. 2010; Sutton et al. 2012, 2013).

To a great surprise, Bachetti et al. (2014) discovered 1.37 second pulsations and 2.5 day orbital modulations in a ULX located in M82, NuSTAR J095551+6940.8, - a clear signature of an accretion on a rotating neutron star. The companion has a minimal mass of $5.2M_{\odot}$, classifying the system as a high mass X -ray binary. This implies isotropic luminosity of the order of a hundred of Eddington luminosities. (Medvedev & Poutanen 2013, discussed a possibility that ULXs are young, rotationally powered pulsars; slow spin and detection of the companion in NuSTAR J095551+6940.8 are evidence of accretion- not rotation-powered source.)

Super-Eddington accretion cannot proceed isotropically (Sec. 2). Slightly super-Eddington accretion rates (Shaviv 1998; Dotan & Shaviv 2011) are possible due to instabilities in the accretion flow. In most cases considered so far, super-Eddington accretion rates were discussed for accretion onto black holes, not neutron stars (see, though Arons 1992). Applicability of these models to the case of neutron star deserves separate investigations. Here we discuss a possible scenario of a geometrically thin, but highly optically thick accretion disk that brings the matter on the surface of a neutron star while producing highly super-Eddington luminosity outside of the disk.

2. Accretion properties

Let us first apply the standard model of accretion in binary systems (*e.g.*, Pringle & Rees 1972; Frank et al. 2002) to the observations of NuSTAR J095551+6940.8. As we will see, the standard accretion model will fail to explain the system's properties, but a simple modified theory succeeds. (In what follows we assume a neutron star with a mass of $1.4M_{\odot}$.)

The system's luminosity is produced due to the accretion of matter on the surface of a neutron star. The X -ray luminosity L_X then is a fraction of the accretion power (Basko & Sunyaev 1976)

$$\begin{aligned} L_{acc} &= \dot{M}GM/R_{NS} \\ L_X &= \eta L_{acc}, \quad \eta \sim 0.5 \end{aligned} \tag{1}$$

(some of the accreted energy is radiated in neutrinos and some is conducted into the neutron star). Parametrizing the luminosity to the Eddington luminosity, $l_E = L_X/L_E$, $L_E = 4\pi cGM_{NS}/\sigma_T$, and the mass accretion rate to the Eddington rate, $\dot{M}_E = L_E/c^2$,

$$\begin{aligned} m_E &= \dot{M}/\dot{M}_E \\ r_G &= 2GM/c^2, \end{aligned} \tag{2}$$

the required m_E is

$$m_E = 2 \frac{l_E}{\eta} \frac{R_{NS}}{r_G} \approx 300 \left(\frac{l_E}{30} \right) \left(\frac{\eta}{0.5} \right), \quad (3)$$

(in physical unites this corresponds to the accretion rate of $5.6 \times 10^{19} \text{ g s}^{-1} = 8.5 \times 10^{-7} M_\odot \text{ yr}^{-1}$). Though the observed luminosity is close to $l_E \sim 100$, the emission is anisotropic; hence we parametrize it to the value of $30L_{Edd}$.

For a given accretion rate m_E the optical depth to Thompson scattering is

$$\tau_T = n\sigma_T r = m_E \sqrt{\frac{r_G}{2r}} \quad (4)$$

Thus, isotropic flow would become optically thick at

$$r \sim 2m_E^2 r_G \approx 7 \times 10^{10} \text{ cm} \quad (5)$$

This would preclude observations of the rotational modulations from the neutron star - the accretion must be anisotropic.

For a given accretion rate the (magnetospheric) Alfvén radius is located at

$$r_A = \left(\frac{B_{NS}^4 R_{NS}^{12}}{2GM\dot{M}^2} \right)^{1/7}. \quad (6)$$

(anisotropy of the accretion flow will reduce r_A , but not substantially due to steep dependence of dipolar magnetic pressure on radius.) It is expected that a system quickly approaches equilibrium, when the Alfvén radius equals the corotation radius r_c

$$r_c = (GM/\Omega^2)^{1/3} = 2 \times 10^8 \text{ cm} \quad (7)$$

The requires surface magnetic field is

$$B_{NS} = 1.4 \times 10^{13} \left(\frac{l_E}{30} \right)^{1/2} \left(\frac{\eta}{0.5} \right)^{-1/2} \text{ G} \quad (8)$$

This value is somewhat larger than the conventional neutron stars' magnetic field of 10^{12} G, but still within the range of population of young pulsars; this is also smaller than the magnetars' magnetic fields (Thompson & Duncan 1995).

All of the above relations follow from the standard accretion models (*e.g.*, Frank et al. 2002). On the other hand, a simple torque prescription for the accreting matter is inconsistent with the above, as already pointed out by Bachetti et al. (2014). The accretion torque due to the material accreting from the disk is

$$I\dot{\Omega} = T_d = \dot{M} \sqrt{GM r_d}, \quad (9)$$

where r_d is the distance at which torque is applied (usually the inner edge of the disk located at r_A). Together with the measured spin-up rate $\dot{P} = -2 \times 10^{-10}$ this implies

$$r_d m_E^2 = 5 \times 10^{10} \text{cm} \quad (10)$$

For $m_E \sim 300$ this would give for r_d an extremely small value, less than the neutron star radius. Thus the accretion torque is more than 10 times smaller than expected from simple accretion models.

Which of the above model assumptions could have gone wrong? The first suspect is the assumption of the corotation at the inner edge of the disk, $r_A = r_c$, *e.g.*, due to the intermittent accretion rate (Bachetti et al. 2014). We disfavor this possibility, since, first, with a period of 1.37 seconds and the spin-up rate of $\dot{P} = -2 \times 10^{-10}$, the time to reach equilibrium, where the inner edge of the disk coincides with corotation, is very short, about 200 hundred years. It is then unlikely that we are observing the source at a special time, with substantially different \dot{M} . Second, from Eqs. (7,10) it follows that the instantaneous accretion rate should be $\langle m_E \rangle \approx 16$, which is nearly 20 times lower than the instantaneous accretion rate (3). Higher accretion efficiency, as well as more anisotropic emission from the accretion column may further alleviate the problem somewhat; still the instantaneous accretion rate must be more than an order of magnitude higher than the average one.

As a resolution of the contradiction between the required high accretion rate and relatively small spin-up rate, we suggest that the torque on the star is reduced with respect to pure disk torque T_d , Eq. (9), due to the penetration of the magnetic field from the star into the accretion disk beyond the corotation, as discussed by Lamb et al. (1973); Arons & Lea (1976); Ghosh & Lamb (1979a,b); Campbell (1987); Aly & Kuijpers (1990); Bardou & Heyvaerts (1996); Agapitou & Papaloizou (2000); Rappaport et al. (2004); Kluźniak & Rappaport (2007); Naso et al. (2013); Wang (1987). Magnetic field lines from the central star penetrate the disk beyond corotation, where they are dragged by the accreting flow with the angular velocity smaller than the angular velocity of the star, thus removing the angular momentum from the star. The additional angular momentum thus deposited in the disk is removed through the conventional disk turbulence.

The exact value of the magnetic torque depends on how the magnetic field of the star penetrates the disk and how it "slides" through the accretion disk - both depend on the assumed prescription for turbulence-generated disk resistivity (Agapitou & Papaloizou 2000). Assuming effective penetration of the magnetic field into the accretion disk, so that on the disk the toroidal magnetic field is $B_\phi \approx -B_z(1 - \omega_K/\Omega_{NS})$ with the poloidal component approximately given by the dipolar field, the magnetic torque is

$$T_M \approx -C(B_{NS}^2 R_{NS}^6)/r_A^3, \quad (11)$$

where C is some diminutional constant (*e.g.*, Rappaport et al. 2004). Since our estimate of the expected torque is nearly two orders of magnitude larger than given by Eq. (9), the matter torque is nearly compensated by the magnetic torque. Equating $T_d = T_M$ gives the magnetospheric radius

$$r_t = C^{2/7} r_A \quad (12)$$

Thus, the parameter C should be of the order of unity (though it enters with a small power) to accommodate large accretion rate, large corotation radius and yet a small torque. For the toroidal magnetic field scaling mentioned above, $C = 1/21$, $C^{2/7} = 0.42$, consistent with the requirement that the magnetospheric radius should be somewhat smaller than the corotation radius for accretion to occur.

Dragging of magnetic fields lines through the disk due to turbulent resistivity (as opposed to twisting due to the frozen-in condition) is valid as long as the difference between the relative linear velocity of solid body rotation of the field lines and of the Keplerian motion in the disk is smaller than the sound speed in the disk, $(\Omega - \sqrt{GM_{NS}/r^3})r \leq c_s$. For a standard α -disk (Shakura & Sunyaev 1973; Frank et al. 2002) this is satisfied only in a fairly narrow region near r_A - the sound speed at r_A is fairly small, $c_s \approx 10^7 \text{ cm s}^{-1}$. Thus, the fastness parameter (Ghosh & Lamb 1979b) must be close to unity. (N.B.: the light cylinder radius $r_{LC} = 6 \times 10^9 \text{ cm}$ is only 30 times larger than the Alfvén radius).

3. Accretion disk and the curtain

3.1. Disk thickness

Using the standard accretion disk theory (Shakura & Sunyaev 1973; Frank et al. 2002) and the parameters estimated above, we can find the thickness of the accretion disk at the Alfvén radius,

$$H \approx 6 \times 10^6 \alpha^{-1/10} \text{ cm}, \quad (13)$$

where α is the disk viscosity parameter (Shakura & Sunyaev 1973). The disk is slightly radiatively dominated, with the ratio of radiation pressure to particle pressure ~ 3 . Thus, the accretion disk is thin, $\Delta \equiv H/r \sim 0.03$. The internal temperature at the inner edge is $T \sim 3 \times 10^6 \text{ K} = 0.3 \text{ keV}$.

In what follows we assume that the relative disk thickness Δ at the Alfvén radius (and not the radial extent of the magnetic field penetration into the disk) determines the relative thickness of the accretion curtain. As a consequence, at a radius $r < r_A$ the cross-section of the accretion curtain is $S = \Delta r r_A$.

3.2. Effects of radiative pressure in the corona

For super-Eddington accretion to occur, most of the infalling material should be shielded from the radiation. Next, we discuss possible geometry of the accretion flow that may allow such process to occur.

Presence of the magnetic field penetrating the disk is the key. Large radiative flux exerts a radial force on the falling matter, mostly on electrons. In magnetic field the radiation pressure leads, first, to drift of electrons in a direction perpendicular both to the force and the local magnetic field, and, second, a force parallel to the magnetic field that in the accretion geometry *pushes the accreting matter towards the magnetic equator*, see Fig. 1. Thus, high central luminosities, combined with penetration of the disk by the stellar magnetic field confine the accreting matter to a narrow accretion curtain.

If a force \mathbf{F} acts on a particle in magnetic field, the particle experiences a drift with velocity

$$\mathbf{v}_d = \frac{\mathbf{F} \times \mathbf{B}}{eB^2}. \quad (14)$$

For an electron in dipolar magnetic field acted upon by the radiative force from a source of luminosity L at the origin this drift velocity equals

$$v_d \approx \frac{\sigma_T L r}{e B_{NS} R_{NS}^3} \quad (15)$$

(at the Alfvén radius (6) this evaluates to several centimeters per second.) The corresponding velocity of the ions is smaller due to smaller scattering cross-section. This drifts of electrons produces a current that modifies the magnetic field, but since we do not expect substantial amount of matter directly exposed to radiation (most is screened), such a modification of the magnetic field is expected to be small.

The kinetic pressure at the edge of the disk, $P \sim 10^{12} \alpha^{-9/10}$ (in cgs units) (Shakura & Sunyaev 1973; Frank et al. 2002) slightly dominates over the radiation pressure, $P_{\text{rad}} \sim L/(4\pi r_A^2 c) \sim 3 \times 10^{11}$. Importantly, due to effective penetration of the stellar magnetic field into the disk this radiative force does not lead to the formation of the wind. For super-Eddington luminosities the parallel repulsive force on the electron is larger than the parallel component of the attractive gravitational force on the ions. Electrons and ions are strongly coupled electrostatically. Both electrons and ions move away from the star, but their motion is restricted by magnetic field. As a results, plasma is radiatively pushed from the corona towards the disk. Thus, the penetrating magnetic field leads to effective “magneto-radiative” compression of the disk. We conclude that super-Eddington luminosity through the disk corona does not lead to the formation of the coronal wind: it keeps the matter in a disk.

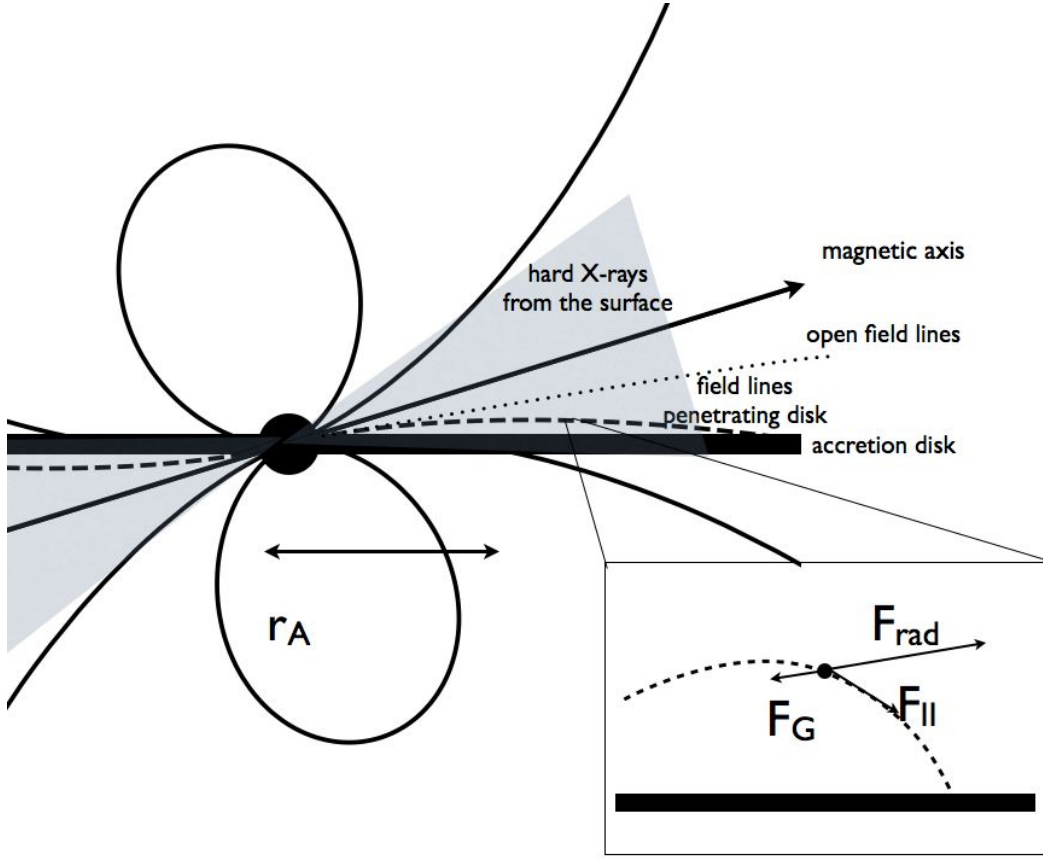


Fig. 1.— Geometry of the accretion curtain. The pulsar is nearly orthogonal rotator with inclination $\geq 80^\circ$ (rotation axis is vertical, perpendicular to the disk surface). Accretion disk is stopped at the Alfvén radius r_A by the magnetospheric magnetic field, and accretes along highly inclined field lines close to the magnetic pole. Accretion occurs in a narrow accretion curtain that cuts through the polar cap. Strong radiative force F_{rad} of the surface hard X -ray emission is higher than the gravity force F_G ; the resulting force parallel to the magnetic field F_{\parallel} clears most of the accretion column and pushes the particles along magnetic fields lines toward the accreting curtain (see the insert). Magnetic fields penetrate the disk beyond the Alfvén radius (dashed field lines); at larger distances the stellar magnetic field lines remain open (dotted line).

The concepts of the well-defined magnetospheric radius and of the effective field penetration into the accretion disk are somewhat contradictory. The assumption is that subdominant fields are turbulently mixed into the disk, while dominant magnetic fields within the Alfvén radius truncate the disk.

(As discussed above, only the parts of the disk close to the inner edge are penetrated by the stellar magnetic fields and are thus subject to “magneto-radiative” compression. Parts of the disk well beyond the Alfvén radius have open field lines and might be subject to radiative ablation. High apparent accretion rates indicate that such ablation is not important, probably due to high incidence angle of the incoming radiation, as well as possible shadowing by the inner edge of the disk).

Note also, that at the magnetospheric radius r_A the radiation pressure, $p_{\text{rad}} \approx L/(4\pi r_A^2 c)$ approximately equals the magnetic field pressure $B(r_A)^2/(8\pi)$. Thus, the coronal field lines can be “pushed” into the disk by the radiation by exerting a radial force on stray particles in the corona.

3.3. Optically thick accretion curtain

As discussed above, due to high radiation pressure material cannot flow along the Alfvén surface - it is pushed back towards the disk by the radiation pressure (if the stellar magnetic field did not penetrate the disk, the matter would be blown away). The material can still accrete if the flow from the accretion disk proceeds along a nearly flat, optically thick accretion curtain. For this, the inclination angle between the magnetic and rotation axis should be large, $\geq 80^\circ$, Fig. 1. In this case accretion proceeds along a nearly flat surface, which cuts across the polar column (see §3.5). The optical depth across the accretion curtain at radius r is

$$\tau_T \approx \Delta n \sigma_T r \approx 4\pi m_E \frac{\sqrt{GM_{NS}r}}{cr_A} \quad (16)$$

where we estimated the density $n = \dot{M}/(m_p S v)$, cross-section of the column $S = \Delta r r_A$ and infall velocity as a free-fall $v \sim \sqrt{GM_{NS}/r}$. At the Alfvén radius $\tau_T \approx 120 \gg 1$.

High optical depth across the accretion curtain and the inner disk has several important implications. First, a large fraction of the hard X -ray surface emission is absorbed by the accretion disk and is then re-radiated. The expected disk luminosity

$$L \approx \pi r_A^2 \sigma_{SB} T^4 \approx 10^{41} T_{\text{keV}}^4 \text{erg s}^{-1} \quad (17)$$

matches the observed luminosity Bachetti et al. (2014).

Second, high optical depth across the accretion curtain implies that accreting matter is mostly shielded from the photons and falls directly onto the surface with nearly free-fall velocity (see also §3.4). Finally, high optical depth across the accretion curtain implies that plasma cannot efficiently cool via sideways emission.

3.4. The last few radii

At sufficiently small distance from the star the magnetic suppression of the scattering cross-section may become important for the X -mode (polarized in a direction perpendicular to the local magnetic field and the direction of propagation), $\sigma_X \approx \sigma_T(\omega/\omega_B)^2$, where ω is the photon frequency. (For nearly radial magnetic field near the magnetic pole most of the surface emission propagates nearly along the field, corresponding to X -mode polarization.) For the given estimate of the magnetic field (8) we can find the distance $r_{Edd,X}$ where the radiative force on electrons $\sim L/(4\pi r^2 c)\sigma_T(k_B T/\hbar\omega_B)^2$ becomes comparable to the gravitational pull on ions:

$$r_{Edd,X} = \sqrt{2\pi}^{1/6} \frac{(e\hbar)^{1/3}(GM)^{5/18}m_p^{1/6}R_{NS}^{1/6}}{(\eta\sigma_T)^{1/6}(k_B T)^{1/3}c^{1/6}m_e^{1/3}\Omega^{7/18}} \approx 1.5R_{NS}T_{NS,10\text{keV}}^{-1/3} \quad (18)$$

Thus, material that gets to $r < r_{Edd,X}$ is freely accreted onto the neutron star.

It is not clear to what temperatures the surface of the neutron star will be heated in such high accretion rates. The direct heating models (*e.g.*, Basko & Syunyaev 1975; Harding et al. 1984) assume low accretion rates, $m_E \ll 1$ and smaller magnetic fields. If for $m_E \gg 1$ the effective surface temperature is tens of keV, then for such high surface temperatures magnetic suppression of the scattering cross-section is not important at the surface magnetic field (8). In addition, resonant scattering may start to become important for surface temperatures ≥ 100 keV. Thus, if the surface temperatures reaches ≥ 50 keV the accretion curtain is mostly shielded from the surface emission all the way to the surface.

On the other hand, for smaller surface temperatures the accreting layer where $\tau \sim 1$ experiences a repulsive radiative force (the thickness of the layer within which the optical depth becomes of the order of unity is $\delta r \approx \sqrt{2}r_{Edd,X}^{3/2}/(\sqrt{r_G}m_E) \approx 2 \times 10^4$ cm). Can this force stop the flow? Comparing the radiation pressure $p_{\text{rad}} \sim L/(4\pi r^2 c)$ with the ram pressure of the infalling material $p_{\text{ram}} = \rho v^2$, $\rho = \dot{M}/(Sv)$, we find that the radiative pressure is smaller than the ram pressure,

$$\frac{p_{\text{rad}}}{p_{\text{ram}}} = \frac{\Delta\eta}{4\sqrt{2}\pi} \sqrt{\frac{r_G}{r}} \frac{r_A}{R_{NS}} \approx 0.1 \quad (19)$$

where the numerical estimate is given for $r = R_{NS}$. Thus, the radiation pressure cannot stop the flow confined to a narrow accretion curtain. (For isotropic flow p_{rad} equals p_{ram} at $r/R_{NS} \approx (2/\eta^2)R_{NS}/r_G \approx 20$.)

(If the radiative pressure was high enough to stop the accreting flow, the highly shielded free-falling accretion flow would lead to accumulation of matter at the $\tau \sim 1$ surface. As a result, the flow would still proceed through formation of narrow accretion channels, somewhat similar to what has been discussed by Shaviv (1998); Dotan & Shaviv (2011); the photon bubble instability can also be important (Arons 1992).)

It is not clear if the accreting matter should pass through a shock near the surface (*e.g.*, Koldoba et al. 2002; Romanova et al. 2002) or can bombard the surface directly (Braun & Yahel 1984; Karino et al. 2008). If the shock forms, an optically thin accretion column at $r \leq 2R_{NS}$ might be an important source of high energy emission.

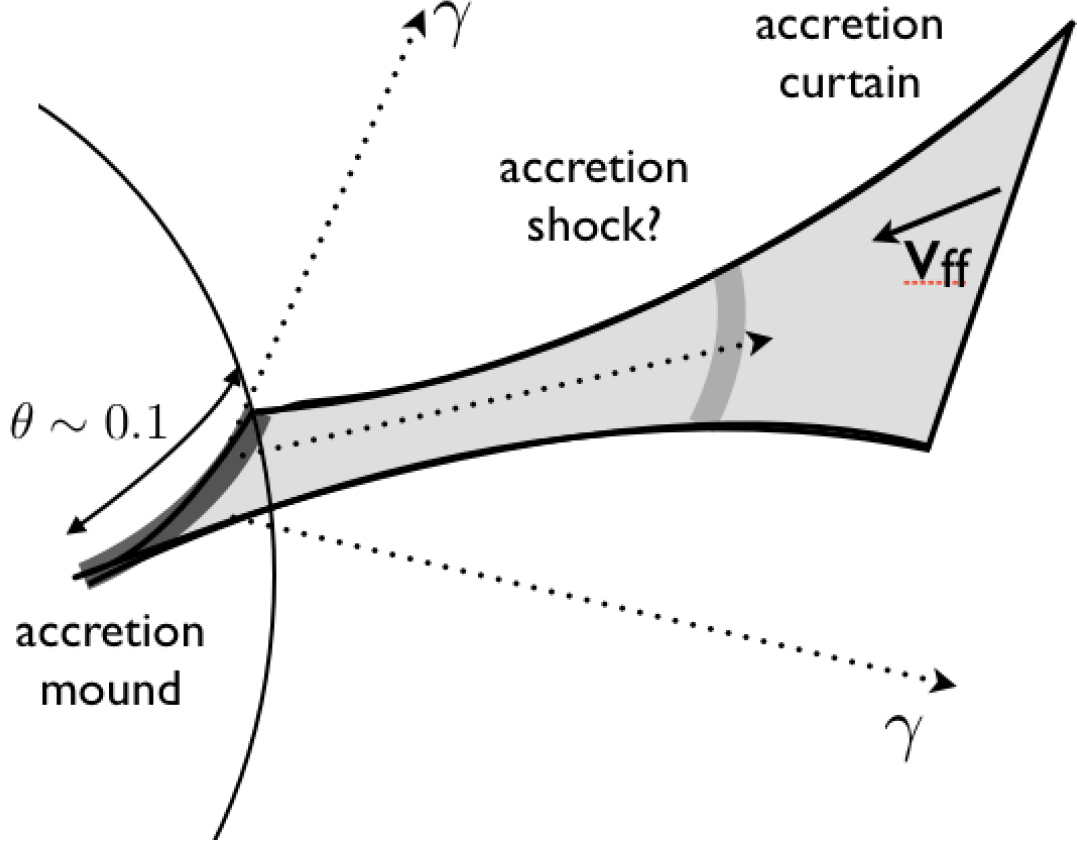


Fig. 2.— Structure of the accretion curtain close to the neutron star. The curtain is nearly flat. The flow is in a free-fall. Below ~ 2 neutron star radii the magnetic suppression of the scattering cross-section allows material to accrete. An accretion shock may form in the flow, even though at higher radii the radiation pressure is still smaller than the ram pressure of the accreting material. The accretion mound is nearly strait narrow region occupying ~ 0.1 in polar angle. Most of the hard X-ray radiation emitted by the mound (dotted lines) escapes the accretion flow.

3.5. The shape of the accretion curtain

Consider magnetic dipole inclined by angle θ_d with respect to the accretion disk’s normal. Let’s assume that the inner edge of the disk corresponds to a fixed magnetic pressure p_A . Also, assume that the disk occupies the plane $\theta = \pi/2$. (It is expected that any corrugation of the disk are smoothed out by the radiation pressure from the central star.) The inner edge of the disk maps along the dipolar magnetic field lines on the points on the surface given by

$$\begin{aligned}\sin \theta &= \frac{\sqrt{1 - \sin^2 \theta_d \cos^2 \phi} \sin \theta_0}{(1 + 3 \sin^2 \theta_d \cos^2 \phi)^{1/12}}, \quad -\pi/2 < \phi < \pi/2 \\ \sin \theta_0 &= (p_A/B_{NS}^2)^{1/12} = \sqrt{R_{NS}/r_A} \ll 1\end{aligned}\tag{20}$$

For aligned case, $\theta_d = 0$, Eq. (20) gives a circle $\theta = \theta_0$, while for orthogonal case, $\theta_d = \pi/2$, Eq. (20) gives a line passing through the magnetic pole and expanding to $\theta = \theta_0$, $\phi = \pm\pi$. In case of nearly orthogonal rotator the accretion will proceed along a nearly flat “curtain”, that extends by $\Delta\theta = \pm\sqrt{R_{NS}/r_A} \sim 0.1$ in the azimuthal direction. For the angle between the rotational and magnetic axes $\theta_d \approx 80^\circ$ the curvature of this accretion curtain near the star is small, $\sim (\pi/2 - \theta_d)/2^{7/6} \sim 0.07$, Fig. 2

4. Discussion

In this Letter we discuss the accretion properties of the highly super-Eddington source ULX NuSTAR J095551+6940.8. We suggest that accretion in this source proceeds in a new regime, which was not studied previously - through an optically thick accretion curtain. Previous studies of slightly super-Eddington accretion flows (Basko & Sunyaev 1976) concluded that most of the accretion energy is radiated through the walls of the accretion column. Highly super-Eddington accretion flows require that most of the accreting matter be shielded from the photons, and, thus, cannot radiate efficiently during the infall.

Several factors are needed for the such super-Eddington luminosity. First, the externally-supplied accretion rate should be sufficiently high. This most likely requires a Roche lobe overflow with either a He-rich companion or a main sequence H-rich companion (Kalogera *et al.* in prep.). Second, the accretion stage should proceed in a new, previously unexplored regime of optically thick accretion curtain. A needed requirement is that the accreting neutron star should be a nearly orthogonal rotator (otherwise an accretion flow along the Alfvén surface need to cover a large solid angle is thus will be exposed to large radiation pressure).

We argue that large accretion rates combined with a small spin-up rate implies that the magnetic field of the neutron star effectively penetrates the accretion disk, removing most

of the angular momentum transferred to the star by the falling material. The suppressions of the torque on the neutron star due to magnetic field penetration has been previously discussed by Ghosh & Lamb (1979b); Campbell (1987); Agapitou & Papaloizou (2000); Wang (1987). This conclusion might have important implications for the estimates of the magnetic fields in the accreting neutron stars. For example, the surface magnetic field found using the luminosity and the assumption of corotation, Eq. (8), is nearly four times larger than would have been inferred using the spin-up rate and the assumption of corotation,

$$B'_{NS} = \frac{\sqrt{2GM_{NS}I_{NS}\dot{\Omega}}}{R_{NS}^3\Omega} = 3 \times 10^{12} \text{ G} \quad (21)$$

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REFERENCES

- Agapitou, V. & Papaloizou, J. C. B. 2000, MNRAS, 317, 273
- Aly, J. J. & Kuipers, J. 1990, A&A, 227, 473
- Arons, J. 1992, ApJ, 388, 561
- Arons, J. & Lea, S. M. 1976, ApJ, 207, 914
- Bachetti, M., Harrison, F. A., Walton, D. J., Grefenstette, B. W., Chakrabarty, D., Fürst, F., Barret, D., Beloborodov, A., Boggs, S. E., Christensen, F. E., Craig, W. W., Fabian, A. C., Hailey, C. J., Hornschemeier, A., Kaspi, V., Kulkarni, S. R., Maccarone, T., Miller, J. M., Rana, V., Stern, D., Tendulkar, S. P., Tomsick, J., Webb, N. A., & Zhang, W. W. 2014, Nature, 514, 202
- Bardou, A. & Heyvaerts, J. 1996, A&A, 307, 1009
- Basko, M. M. & Sunyaev, R. A. 1976, MNRAS, 175, 395
- Basko, M. M. & Syunyaev, R. A. 1975, Soviet Journal of Experimental and Theoretical Physics, 41, 52
- Braun, A. & Yahel, R. Z. 1984, ApJ, 278, 349
- Campbell, C. G. 1987, MNRAS, 229, 405
- Dotan, C. & Shaviv, N. J. 2011, MNRAS, 413, 1623

- Fabbiano, G. & White, N. E. Compact stellar X-ray sources in normal galaxies, ed. W. H. G. Lewin & M. van der Klis, 475–506
- Feng, H. & Soria, R. 2011, *New A Rev.*, 55, 166
- Frank, J., King, A., & Raine, D. J. 2002, *Accretion Power in Astrophysics: Third Edition*
- Ghosh, P. & Lamb, F. K. 1979a, *ApJ*, 232, 259
- . 1979b, *ApJ*, 234, 296
- Harding, A. K., Kirk, J. G., Galloway, D. J., & Meszaros, P. 1984, *ApJ*, 278, 369
- Kaaret, P., Corbel, S., Prestwich, A. H., & Zezas, A. 2003, *Science*, 299, 365
- Karino, S., Kino, M., & Miller, J. C. 2008, *Progress of Theoretical Physics*, 119, 739
- Kluźniak, W. & Rappaport, S. 2007, *ApJ*, 671, 1990
- Koldoba, A. V., Lovelace, R. V. E., Ustyugova, G. V., & Romanova, M. M. 2002, *AJ*, 123, 2019
- Lamb, F. K., Pethick, C. J., & Pines, D. 1973, *ApJ*, 184, 271
- Mapelli, M., Ripamonti, E., Zampieri, L., Colpi, M., & Bressan, A. 2010, *MNRAS*, 408, 234
- Medvedev, A. S. & Poutanen, J. 2013, *MNRAS*, 431, 2690
- Naso, L., Kluźniak, W., & Miller, J. C. 2013, *MNRAS*, 435, 2633
- Pringle, J. E. & Rees, M. J. 1972, *A&A*, 21, 1
- Rappaport, S. A., Fregeau, J. M., & Spruit, H. 2004, *ApJ*, 606, 436
- Roberts, T. P. 2007, *Ap&SS*, 311, 203
- Romanova, M. M., Ustyugova, G. V., Koldoba, A. V., & Lovelace, R. V. E. 2002, *ApJ*, 578, 420
- Shakura, N. I. & Sunyaev, R. A. 1973, *A&A*, 24, 337
- Shaviv, N. J. 1998, *ApJ*, 494, L193
- Sutton, A. D., Roberts, T. P., & Middleton, M. J. 2013, *MNRAS*, 435, 1758

Sutton, A. D., Roberts, T. P., Walton, D. J., Gladstone, J. C., & Scott, A. E. 2012, MNRAS, 423, 1154

Thompson, C. & Duncan, R. C. 1995, MNRAS, 275, 255

Wang, Y.-M. 1987, A&A, 183, 257